

# Split NMSSM with electroweak baryogenesis

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## Abstract

In light of the Higgs boson discovery we reconsider generation of the baryon asymmetry in the non-minimal split Supersymmetry model with an additional singlet superfield in the Higgs sector. We find that successful baryogenesis during the first order electroweak phase transition is possible within phenomenologically viable part of the model parameter space. We discuss several phenomenological consequences of this scenario, namely, predictions for the electric dipole moments of electron and neutron and collider signatures of light charginos and neutralinos.

## 1 Introduction

Any phenomenologically viable particle physics model should explain the observed asymmetry between matter and antimatter in the Universe. The analysis of the anisotropy and polarization of the cosmic microwave background provided by WMAP collaboration gives the following baryon-to-photon ratio [1]

$$\frac{n_B}{n_\gamma} = (6.19 \pm 0.14) \times 10^{-10}. \quad (1)$$

To generate the baryon asymmetry of the Universe, three Sakharov's conditions should be satisfied [2]: (i) baryon number violation, (ii)  $C$ - and  $CP$ -violation and (iii) departure from thermal equilibrium. The latter condition can be realized, in particular, during the strong first order electroweak phase transition (EWPT) which proceeds via nucleation and expansion of bubbles of new phase in the hot plasma of the early Universe (for a recent discussion see, e.g., Refs. [3, 4]). The baryon number violation during the EWPT happens due to sphaleron processes in symmetric phase, while the  $CP$ -violation is induced by the interaction of particles in plasma with the bubble walls.

In the Standard Model of particle physics (SM) the Sakharov's conditions are only partly fulfilled. In particular, baryon number is violated via electroweak sphaleron transitions at high temperatures. At the same time, the electroweak transition in the SM is not the first order phase transition, hence no sufficient departure from thermal equilibrium. And the contribution of  $CP$ -violating CKM phases is too small in any case to provide

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(1). Finally, the electroweak sphalerons in the broken phase are too fast and would wash out any baryon asymmetry generated during the EWPT [5, 6]. Therefore, electroweak baryogenesis is only possible in SM extensions. These models should contain additional sources of  $CP$ -violation. Moreover, if the baryon asymmetry emerges at the electroweak scale, there should be a mechanism making the EWPT to be the strongly first order. A lot of scenarios for baryogenesis during the EWPT have been proposed and studied, see e.g. Refs. [7, 8, 9, 10, 11, 12, 13, 14, 15].

The Minimal Supersymmetric Standard Model (MSSM) is one of the most elegant ways to extend the SM framework. In particular, the quadratic divergences cancellation and the gauge couplings unification are the major reasons for the interest in supersymmetric models. Moreover, the lightest neutralino is a natural dark matter candidate in the MSSM [16, 17]. In general, however, the Higgs boson discovery [18, 19], and non-observation of superpartners at the LHC shrinks severely the region of MSSM parameter space. For instance, squarks and gluinos have been searched for at the LHC [20, 21], and the lower bounds on their masses have been set at the level of 1-2 TeV.

An attractive MSSM extension with splitted superpartner spectrum (split MSSM) has been proposed in Refs. [22, 23]. The squarks and sleptons in these scenarios are very heavy, while neutralinos and charginos remain light. Nevertheless, the main advantages of SUSY, i.e. the gauge coupling unification and existence of dark matter candidate, remain intact in this class of models. Remarkably, the absence of FCNC processes [24] is naturally understood within this setup. Unfortunately, the electroweak baryogenesis can not be realized in minimal version of the split SUSY. This can be cured by introducing a gauge singlet superfield to the Higgs sector of the split MSSM [25]. The main features of this split Next-to-Minimal Supersymmetric extension of the Standard Model, *split NMSSM*, are the following. There are two energy scales in the split NMSSM, electroweak  $M_{EW} \sim 100 \text{ GeV}$  and splitting scale  $M_S \gg M_{EW}$ . At  $M_{EW}$  scale, the spectrum of split NMSSM contains the SM particles, one Higgs doublet  $H$ , the higgsino components  $\tilde{H}_{u,d}$ , winos  $\tilde{W}$ , bino  $\tilde{B}$ , and in addition a singlet complex scalar field  $N$  and its superpartner singlino  $\tilde{n}$ . The sleptons, squarks and four out of seven scalar degrees of freedom in the Higgs sector have masses of order the splitting scale  $M_S$ . Hence, these particles are decoupled from the spectrum at low energies  $E < M_S$ . At the same time, interactions of the scalar components of the singlet  $N$  with the Higgs boson are described at  $M_{EW}$  by a generic potential, which includes trilinear terms. These couplings are capable of strengthening the first order EWPT. In the present paper, we review this scenario in view of the latest experimental results, in particular, the Higgs boson discovery.

This paper is organized as follows. In Section 2 we discuss the structure of split NMSSM. In Section 3, we explore the phenomenologically allowed region of the model parameters consistent with the Higgs boson of mass  $m_H \simeq 125 \text{ GeV}$ . In Sections 4 and 5 we study the strong first order EWPT and the baryon asymmetry of the Universe, respectively, for the relevant split NMSSM parameter space. In Section 6 we perform an analysis of the electron and neutron EDMs. There we also discuss the spectra of charginos and neutralinos, which can be probed at the LHC experiments. In Appendix A we calculate one-loop renormalization group (RG) corrections to the Higgs boson mass,

which are needed to find allowed region of the parameter space in the split NMSSM scenario. In Appendix B the minimization conditions for the split NMSSM effective potential are presented.

## 2 Non-minimal split Supersymmetry

In this Section we discuss the Lagrangian and particle content of the split NMSSM. Above the splitting scale  $M_S$ , the model is described by generic<sup>4</sup> NMSSM superpotential

$$W = \lambda \hat{N} \hat{H}_u \epsilon \hat{H}_d + \frac{1}{3} k \hat{N}^3 + \mu \hat{H}_u \epsilon \hat{H}_d + r \hat{N}, \quad (2)$$

where  $\hat{H}_{u,d}$  are superfields of the Higgs doublets,  $\hat{N}$  is a chiral superfield singlet with respect to  $SU(3)_C \times SU(2)_L \times U(1)_Y$  gauge group,  $\hat{N} = N + \sqrt{2} \theta \tilde{n} + \theta^2 F_N$ , and  $\epsilon$  is antisymmetric  $2 \times 2$  matrix with  $\epsilon_{12} = -\epsilon_{21} = 1$ .

The tree level scalar potential of the non-minimal SUSY model can be written as follows

$$V = V_D + V_F + V_{soft}, \quad (3)$$

where the contribution of  $D$ -terms is the same as that in the MSSM,

$$V_D = \frac{g^2}{8} \left( H_d^\dagger \sigma_a H_d + H_u^\dagger \sigma_a H_u \right)^2 + \frac{g'^2}{8} (|H_d|^2 - |H_u|^2)^2,$$

with  $g$  and  $g'$  being  $SU(2)_L$  and  $U(1)_Y$  gauge couplings, respectively. The contribution of  $F$ -terms derived from superpotential (2) reads

$$V_F = |\lambda H_u \epsilon H_d + k N^2 + r|^2 + |\lambda N + \mu|^2 \left( H_u^\dagger H_u + H_d^\dagger H_d \right).$$

Soft supersymmetry breaking terms are described by the potential

$$V_{soft} = \left( \lambda A_\lambda N H_u \epsilon H_d + \frac{1}{3} k A_k N^3 + \mu B H_u \epsilon H_d + A_r N + \text{h.c.} \right) \quad (4)$$

$$+ m_u^2 H_u^\dagger H_u + m_d^2 H_d^\dagger H_d + m_N^2 |N|^2, \quad (5)$$

where  $A_{\lambda,k}$  and  $m_{u,d,N}$  are the trilinear couplings and the soft masses of scalars, respectively. Components of the Higgs doublets  $H_{u,d}$  and singlet field  $N$  in (4), (5) are defined by

$$H_u = \begin{pmatrix} H_u^+ \\ H_u^0 \end{pmatrix}, \quad H_d = \begin{pmatrix} H_d^0 \\ H_d^- \end{pmatrix}, \quad N = (S + iP)/\sqrt{2}, \quad (6)$$

where  $S$  and  $P$  are the scalar and pseudoscalar parts of the singlet  $N$ , correspondingly. We introduce the following notations:  $\tan \beta \equiv \langle H_u^0 \rangle / \langle H_d^0 \rangle$ ,  $v_S \equiv \langle S \rangle$  and  $v_P \equiv \langle P \rangle$ .

An explicit analysis of the particle spectrum of the model with the potential (3) is performed in Ref. [25]. We nevertheless briefly discuss the particle content of the scalar sector at energies below the splitting scale. There are ten scalar degrees of freedom at

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<sup>4</sup>A quadratic in  $\hat{N}$  term can be eliminated by a field redefinition.

the splitting scale  $M_S$ , coming from (6). It is shown in Ref. [25] that if the soft SUSY breaking parameters  $B\mu$ ,  $m_d^2$  and  $m_u^2$  are of order of the squared splitting scale,  $M_S^2$ , then two charged Higgses, one pseudoscalar and one neutral scalar Higgs bosons are heavy and thus decoupled from the low energy spectrum, while a fine-tuning is required for the mass of the lightest Higgs boson  $H$  and two singlets,  $S, P$  to be at the electroweak scale. Three Goldstone modes are eaten by  $W^\pm$  and  $Z^0$  due to the Higgs mechanism. We emphasize that the particle spectrum in the split NMSSM (as well as in any split SUSY model) below  $M_S$  requires a fine-tuning of the soft dimensionful parameters [25].

Replacing  $H_u \rightarrow H \sin \beta$  and  $H_d \rightarrow \epsilon H^* \cos \beta$  in (3) we obtain at the splitting scale  $M_S$  the effective Lagrangian for the relevant at low energy degrees of freedom in the scalar sector of the model (hereafter we omit the corresponding kinetic terms),

$$\begin{aligned} -\mathcal{L}_V = & \frac{\bar{g}^2}{8} \cos^2 2\beta (H^\dagger H)^2 + |r + kN^2 - \frac{\lambda}{2} \sin 2\beta H^\dagger H|^2 + |\lambda N + \mu|^2 H^\dagger H \\ & + \left( -\frac{\lambda}{2} A_\lambda \sin 2\beta N H^\dagger H - \frac{\mu B}{2} \sin 2\beta H^\dagger H + \frac{1}{3} k A_k N^3 + A_r N + h.c. \right) \\ & + (m_u^2 \sin^2 \beta + m_d^2 \cos^2 \beta) H^\dagger H + m_N^2 |N|^2, \end{aligned} \quad (7)$$

where  $\bar{g}^2 \equiv g^2 + (g')^2$ . The quark-Higgs Yukawa interactions, gaugino couplings and gaugino mass terms are the same as in the minimal split supersymmetry. New part of the Yukawa interactions for Higgsinos  $\tilde{H}_{u,d}$  and singlino field  $\tilde{n}$  is given by

$$-\mathcal{L}_Y = -\lambda N \tilde{H}_u \epsilon \tilde{H}_d - \lambda \sin \beta H^T \epsilon (\tilde{H}_d \tilde{n}) + \lambda \cos \beta (\tilde{n} \tilde{H}_u) H^* - k N \tilde{n} \tilde{n} + h.c. \quad (8)$$

Now we consider the most general scalar Lagrangian at energies below  $M_S$

$$\begin{aligned} -\mathcal{L}_V = & -m^2 H^\dagger H + \frac{\tilde{\lambda}}{2} (H^\dagger H)^2 + i \tilde{A}_1 H^\dagger H (N^* - N) + \tilde{A}_2 H^\dagger H (N + N^*) + 2\kappa_1 |N|^2 H^\dagger H \\ & + \kappa_2 H^\dagger H (N^2 + N^{*2}) + \tilde{m}_N^2 |N|^2 + \lambda_N |N^2|^2 + \frac{1}{3} \tilde{A}_k (N^3 + N^{*3}) + \tilde{A}_r (N + N^*) \\ & + \left( \frac{\tilde{m}^2}{2} N^2 + \frac{1}{2} \tilde{A}_3 N^2 N^* + \xi N^4 + \frac{\eta}{6} N^3 N^* + h.c. \right), \end{aligned} \quad (9)$$

here the quartic couplings  $\tilde{\lambda}$ ,  $\kappa$ ,  $\kappa_1$ ,  $\kappa_2$  and  $\lambda_N$  at the electroweak scale are related via renormalization group equations to  $\bar{g}$ ,  $\lambda$ ,  $k$  and  $\tan \beta$  at the scale  $M_S$ . Comparing scalar potential (9) with (7) one can obtain the matching conditions for these couplings at the splitting scale  $M_S$ :

$$\kappa_1 = \lambda^2, \quad \kappa_2 = -\lambda k \sin \beta \cos \beta, \quad \lambda_N = k^2, \quad \kappa = \lambda, \quad (10)$$

$$\tilde{\lambda} = \frac{\bar{g}^2}{4} \cos^2 2\beta + \frac{\lambda^2}{2} \sin^2 2\beta. \quad (11)$$

We use here the convention  $g_1^2 = (5/3)g'^2$  and  $g_2 = g$  adopted in Grand Unified Theories (GUT). Note that the couplings proportional to  $\xi$  and  $\eta$  in (9) are absent in the effective

Lagrangian at  $M_S$ , but get induced by loop quantum corrections; thus we set the following RG initial condition

$$\xi = \eta = 0 \quad (12)$$

at the splitting scale  $M_S$ . Soft fermion masses and Yukawa interactions below  $M_S$  are described by the Lagrangian

$$\begin{aligned} -\mathcal{L}_Y = & \frac{M_2}{2} \tilde{W}^a \tilde{W}^a + \frac{M_1}{2} \tilde{B} \tilde{B} + (\mu + \kappa N) \tilde{H}_u^T \epsilon \tilde{H}_d - k N \tilde{n} \tilde{n} \\ & + H^\dagger \left( \frac{1}{\sqrt{2}} \tilde{g}_u \sigma^a \tilde{W}^a + \frac{1}{\sqrt{2}} \tilde{g}'_u \tilde{B} - \lambda_u \tilde{n} \right) \tilde{H}_u \\ & + H^T \epsilon \left( -\frac{1}{\sqrt{2}} \tilde{g}_d \sigma^a \tilde{W}^a + \frac{1}{\sqrt{2}} \tilde{g}'_d \tilde{B} - \lambda_d \tilde{n} \right) \tilde{H}_d + h.c., \end{aligned} \quad (13)$$

where  $M_2$  and  $M_1$  are wino and bino soft mass parameters in  $SU(2)_L$  and  $U(1)_Y$  gaugino sectors, respectively. The corresponding matching conditions for Yukawa couplings at the splitting scale  $M_S$  read

$$\lambda_u = \lambda \cos \beta, \quad \lambda_d = -\lambda \sin \beta, \quad (14)$$

$$\tilde{g}_u = g \sin \beta, \quad \tilde{g}_d = g \cos \beta, \quad \tilde{g}'_u = g' \sin \beta, \quad \tilde{g}'_d = g' \cos \beta. \quad (15)$$

Matching equations for the dimensionful couplings in (9) can be found in a similar way. However, for simplicity we take their values directly at electroweak scale rather than solving RG equations for them from  $M_S$  down to electroweak energies. In order to reduce the number of trilinear couplings we assume that Higgs-scalar ( $H-S$ ) and Higgs-pseudoscalar ( $H-P$ ) mixing terms in their squared mass matrix are equal to zero at the EW energy scale. This implies appropriate relations for the trilinear couplings  $\tilde{A}_1$  and  $\tilde{A}_2$ ,

$$\tilde{A}_1 = \sqrt{2}(\kappa_1 - \kappa_2)v_P, \quad \tilde{A}_2 = -\sqrt{2}(\kappa_1 + \kappa_2)v_S. \quad (16)$$

From the very beginning we admit explicit  $CP$ -violation by taking purely imaginary  $\mu$ -term and from largangians (9) and (7) we relate its value through the following matching condition at  $M_S$  scale

$$\text{Im } \mu = \tilde{A}_1 / \lambda \quad (17)$$

neglecting small RG corrections. Let us note that using minimization conditions for the potential (9), soft squared masses  $m^2$ ,  $\tilde{m}^2$  and  $\tilde{m}_N^2$  can be re-expressed via vevs of the scalar fields, i.e.  $v$ ,  $v_S$  and  $v_P$ . For completeness these relations are presented in Appendix B.

With the all above assumptions, we are left with only seven independent dimensionful parameters of the model at the EW scale

$$(v_S, \quad v_P, \quad M_1, \quad M_2, \quad \tilde{A}_k, \quad \tilde{A}_3, \quad \tilde{A}_r). \quad (18)$$

In what follows to get numerical results, for concreteness, we set at the EW scale:

$$\tilde{A}_3 = \tilde{A}_r = 0, \quad \tilde{A}_k = -1.1 \text{ GeV}, \quad (19)$$

while scanning over all the other four parameters. We advertise that the two singlet VEVs  $v_S$  and  $v_P$  play very prominent role in developing the EWPT, which is discussed below in Section 4.

### 3 Predictions for the Higgs boson mass

In this Section we describe the scanning over the set of three dimensionless parameters  $(\tan\beta, \lambda, k)$  fixed at scale  $M_S$  and calculate the mass of the Higgs boson resonance. We outline the region of model parameter space consistent with the SM-like Higgs boson with mass about 125 GeV.

In our procedure we choose dimensionless couplings of the model at the splitting scale and calculate the value of the Higgs boson mass by solving RG equations at next-to-leading order in coupling constants (NLO). We start solving the truncated part of the RG equations from the EW up to the splitting scale for the SM couplings

$$(g', \quad g, \quad g_s, \quad y_t), \quad (20)$$

where  $g_s$  is  $SU(3)_c$  gauge coupling and  $y_t$  denotes the top Yukawa coupling. Initial conditions for RG equations for these couplings at the EW scale are taken as follows [24]

$$\alpha_s(M_Z) = 0.118, \quad M_Z = 91.19 \text{ GeV}, \quad M_W = 80.39 \text{ GeV}, \quad \text{and} \quad y_t(m_t) = 0.95.$$

Next, we use complete set of the RG equations for dimensionless couplings of the split NMSSM

$$(g', \quad g, \quad g_s, \quad y_t, \quad \tilde{\lambda}), \quad (\tilde{g}_{u,d}, \quad \tilde{g}'_{u,d}, \quad \lambda_{u,d}, \quad \kappa, \quad \kappa_{1,2}, \quad k, \quad \lambda_N, \quad \xi, \quad \eta). \quad (21)$$

Corresponding RG equations can be found in Ref. [25]. In order to obtain values of the couplings (21) at low energies, the values of  $\tan\beta$ ,  $\lambda$  and  $k$  are chosen randomly at the splitting scale  $M_S$  from the following perturbative regions

$$-0.6 < k < 0.6, \quad 0 < \lambda < 0.7, \quad 0 < \tan\beta < 30. \quad (22)$$

Then we solve the complete set of the RG equations from  $M_S$  down to the EW scale by using matching condition (10), (12), (14) and (15). This procedure doesn't guarantee the correct value of top Yukawa coupling at low energy  $y_t(M_t)$ . Therefore, we tune  $y_t(M_S)$  to obtain the value of  $y_t(M_t)$  within the error bars (for details see A.3 and Refs. [25, 26]). We include a part of threshold correction to the Higgs quartic coupling at the splitting scale [27] resulted in the following modification

$$\tilde{\lambda} \rightarrow \tilde{\lambda} + \delta\tilde{\lambda}, \quad (23)$$

where  $\delta\tilde{\lambda}$  is a conversion term from  $\overline{\text{DR}}$  to  $\overline{\text{MS}}$  renormalization schemes at  $M_S$ ,

$$\delta\tilde{\lambda} = -\frac{1}{16\pi^2} \left[ \frac{9}{100} g_1^4 + \frac{3}{10} g_1^2 g_2^2 + \left( \frac{3}{4} - \frac{\cos^2 2\beta}{6} \right) g_2^4 + \frac{3}{400} (5g_2^2 + 3g_1^2)^2 \sin^2 4\beta \right]. \quad (24)$$

The remaining part of the threshold correction to  $\tilde{\lambda}$  depends on hierarchy of masses of heavy scalars near the splitting scale and it has not been taken into account. We should keep it in mind when interpreting the results. Next, we calculate the pole mass of the Higgs boson including one-loop threshold corrections at the electroweak scale, see Appendix A

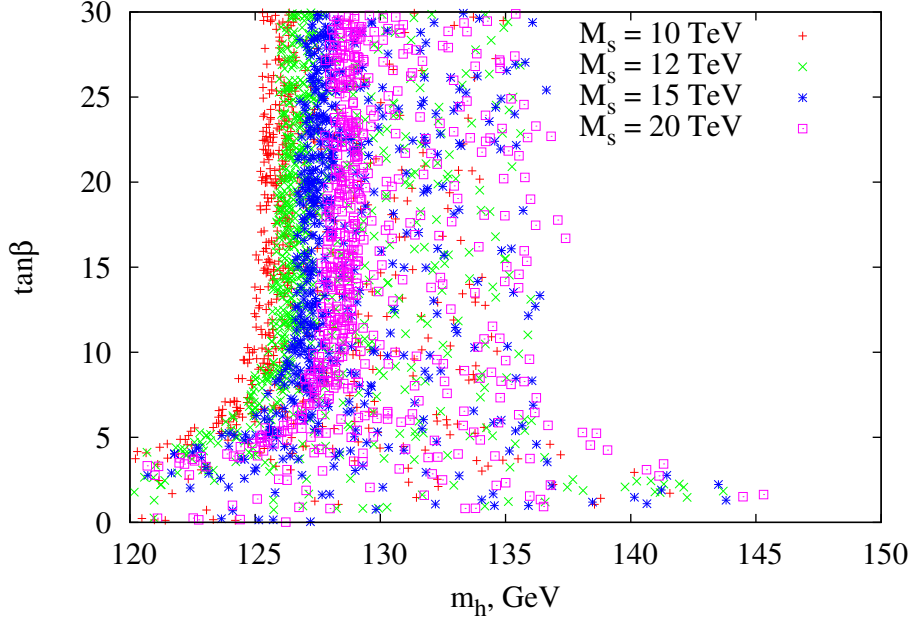


Figure 1: Prediction for the Higgs boson mass  $m_h$  as a function of  $M_S$  and  $\tan\beta$ . We assumed here that the Yukawa top coupling falls within the range  $y_t^{lower} < y_t < y_t^{upper}$ , see the main text for details.

for details. In Fig. 1 we show prediction for the Higgs boson mass obtained with various values of split scale  $M_S$  and  $\tan\beta$ . It follows from Fig. 1 that for most of the models the Higgs mass shifts by several GeVs if one increases the splitting scale  $M_S$  from 10 to 20 TeV for  $\tan\beta > 10$ . The similar behavior was observed in split MSSM [27]. This is attributed to a large quantum correction coming from heavy stops.

Now, we require that the pole mass of the Higgs boson (52) and  $y_t$  at  $\bar{\mu} = M_t$  fall within the following ranges

$$125.3 \text{ GeV} < m_h^{pole} < 125.9 \text{ GeV}, \quad y_t^{lower} < y_t < y_t^{upper}.$$

Here we use the average value  $m_h = 125.6 \pm 0.3 \text{ GeV}$  from CMS [18] and ATLAS [19] combined results (for details see, e.g., Ref. [24] and references therein). Lower and upper limits for  $y_t$  are extracted from Eq. (80) and correspond to  $M_t^{lower} = 172.3 \text{ GeV}$  and  $M_t^{upper} = 174.1 \text{ GeV}$  respectively. In Fig. 2 we show the selected models in  $(\tan\beta, \lambda)$ -plane for the values of the splitting scale  $M_S$  varying from 10 to 20 TeV. One can see that for  $\tan\beta > 5$  parameter  $\lambda$  can take arbitrary values in the allowed perturbative region. For  $\tan\beta \simeq 1$  the allowed region shifts to the maximal values of  $\lambda$  which follows from the matching condition (11). We check that  $\lambda$  is in the perturbative regime up to the GUT scale. In addition, as follows from Fig. 2, the phenomenologically possible values of  $\tan\beta$  grow with decreasing of the splitting scale  $M_S$  for  $\lambda < 0.4$ . This is again related to the balance between the tree-level and loop-induced contributions to the Higgs boson mass. The regions where  $\tan\beta$  is either large ( $\beta \rightarrow \pi/2$ ) or small ( $\beta \rightarrow 0$ ) correspond to the

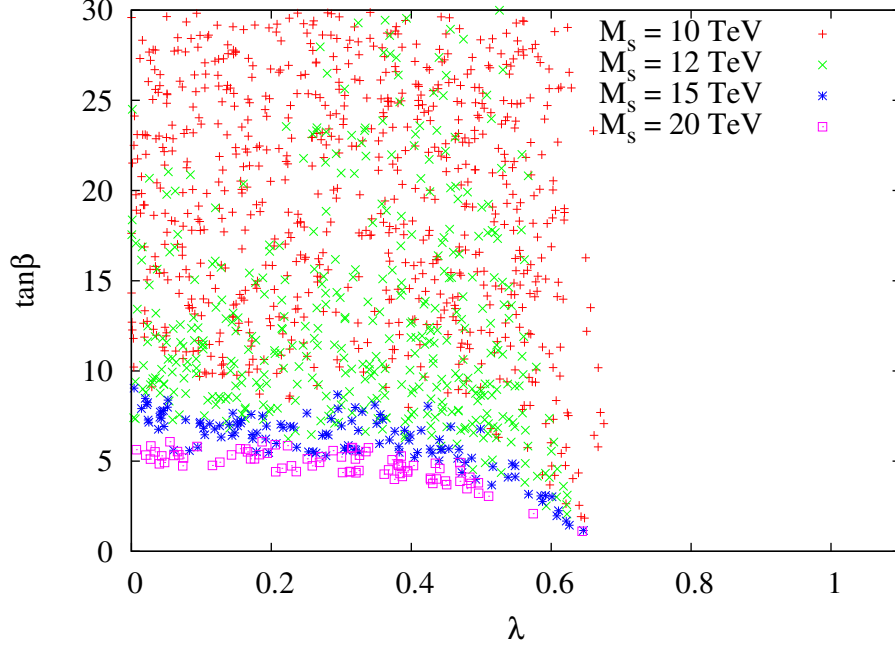


Figure 2: Allowed regions for  $\tan \beta$  and  $\lambda(M_S)$  for various values of the splitting scale  $M_S$ .

	$\tilde{g}_u$	$\tilde{g}_d$	$\tilde{g}'_u$	$\tilde{g}'_d$	$\lambda_u$	$\lambda_d$	$\kappa$	$\kappa_1$	$\kappa_2$	$\lambda_N$
Setup 1	0.650	0.070	0.347	0.037	0.057	-0.513	0.560	0.251	-0.022	0.208
Setup 2	0.649	0.065	0.347	0.034	0.056	-0.560	0.609	0.297	-0.021	0.207

Table 1: Dimensionless couplings at the electroweak scale.

decoupling of the second term in (11). We find that for  $M_S \rightarrow \infty$  the allowed regions for  $\tan \beta$  and  $\lambda$  shrink to  $\tan \beta \rightarrow 1$  and  $\lambda \rightarrow 0$ , respectively.

As it follows from (10) and (11) the coupling  $k$  does not enter the matching condition for  $\tilde{\lambda}$  at  $M_S$  and we find that the value of the Higgs boson mass in the model is almost independent of the coupling constant  $k$  within the perturbative ranges (22).

In what follows, we choose two close benchmark setups for the free parameters

$$\text{Setup 1 : } \quad M_S = 12 \text{ TeV}, \tan \beta = 9.21, \lambda = 0.559, k = -0.5; \quad (25)$$

$$\text{Setup 2 : } \quad M_S = 10 \text{ TeV}, \tan \beta = 10.0, \lambda = 0.611, k = -0.5. \quad (26)$$

The both benchmark models are well inside the allowed regions in Fig. 2. For calculation of the threshold correction the relevant dimensionful parameters are taken to be  $M_2 = 1 \text{ TeV}$ ,  $M_1 = 300 \text{ GeV}$  and  $\text{Im} \mu = 120 \text{ GeV}$ . As it has been found in [25] the resulting baryon asymmetry is directly related to the value of  $\lambda$ . Thus the coupling  $\lambda$  is rather large for both chosen models. The relevant Yukawa and quartic couplings at the electroweak scale,  $\bar{\mu} = M_t = 173.2 \text{ GeV}$ , are presented in Table 1. Below we use these couplings in the



analysis of the strong first order EWPT (Section 4), in the calculation of BAU (Section 5) and to estimate the values of EDMs of the electron and neutron (Section 6).

## 4 Strong first order EWPT

In this Section we revisit the results of Ref. [25] for the strongly first order electroweak phase transition in the split NMSSM within the region of the parameter space favored by the measured value of the Higgs boson mass ( $m_h \simeq 125$  GeV). Let us consider the effective potential at finite temperature  $T$  [28]

$$V_T^{eff} = V_{T=0}^{eff} + V_T^{(1)}. \quad (27)$$

Here  $V_T^{(1)}$  is the thermal contribution given by

$$V_T^{(1)} = \sum_i f_i(m_i, T), \quad (28)$$

where sum goes over all species in the hot plasma (see, e.g., Eq. (85)), and

$$f_i(m_i, T) = (\pm) \frac{T^4}{2\pi^2} \int_0^\infty dx x^2 \ln \left( 1 \mp e^{-\sqrt{x^2 + (m_i/T)^2}} \right), \quad (29)$$

where the upper and lower signs correspond to bosons and fermions respectively.

In order to avoid baryon number washout after the phase transition the condition  $v_c/T_c \gtrsim 1.1$  has to be satisfied [29] (see also recent revised discussion in [30]). Here  $v_c$  is the Higgs VEV at the critical temperature  $T_c$ . We define  $T_c$  as a temperature at which one bubble of the broken phase begins to nucleate within a causal space-time volume of the Universe. The latter is determined by the Hubble parameter  $\mathcal{H}(T)$  as

$$\mathcal{H}^{-4}(T) = (M_{Pl}^*/T^2)^4. \quad (30)$$

The bubble nucleation rate in a unit space-volume has the form

$$\Gamma(T) \simeq (\text{prefactor}) \times T^4 \exp(-S_3/T), \quad (31)$$

where  $S_3 = S_3(T)$  is the free energy of the critical bubble at a given temperature

$$S_3(T) = 4\pi \int_0^\infty dr r^2 \left[ \frac{1}{2} \left( \frac{dh}{dr} \right)^2 + \frac{1}{2} \left( \frac{dS}{dr} \right)^2 + \frac{1}{2} \left( \frac{dP}{dr} \right)^2 + V_T^{eff}(h, S, P) \right]. \quad (32)$$

Here  $h(r)$ ,  $S(r)$  and  $P(r)$  are the radial configurations of the scalar fields, which minimize the functional  $S_3$ . Therefore, the probability that the bubble is nucleated inside a causal volume reads

$$P \sim \Gamma \cdot \mathcal{H}^{-4} \sim \frac{M_{Pl}^{*4}}{T^4} \exp(-S_3/T). \quad (33)$$

The first bubble nucleates when  $P \sim 1$ , which yields a rough estimate for the nucleation criterion  $S_3(T)/T \sim 4 \ln \left( \frac{M_{Pl}^*}{T} \right) \sim 150$ , where  $T$  is a typical temperature of order the electroweak energy scale,  $T \simeq M_{EW}$ . More accurate calculation reveals [31]

$$S_3(T_c)/T_c \simeq 135. \quad (34)$$

We recall that singlet VEVs  $v_S$  and  $v_P$  are the input parameters of our model. The vacuum  $(v, v_S, v_P)$  is the global minimum of the effective potential  $V_{T=0}^{eff}$  in the broken phase (see discussion in Appendix B). At the finite temperature  $T \neq 0$ , this broken minimum is shifted due to the thermal corrections (28)

$$(v, v_S, v_P) \rightarrow (v_c, S_c, P_c). \quad (35)$$

In order to find numerically the profile of the critical bubbles, we use the method described in [32, 33] and later modified in Ref. [25]. The procedure can be summarized as follows. Firstly, for given values of  $v_S$ ,  $v_P$  and  $T$  we find numerically the nearest minima of the effective potential  $V_T^{eff}$  in the symmetric  $(0, S_s, P_s)$  and broken  $(v_c, S_c, P_c)$  phases. For technical reason, we shift the effective potential by a constant to set  $V_T^{eff}(0, S_s, P_s) = 0$ . Secondly, we construct an ansatz for the bubble wall configurations which interpolate between these two minima at the temperature  $T$ . Next, using this ansatz as the first approximation we numerically find the absolute minimum of the functional

$$\mathcal{F}(h, S, P) = 4\pi \int_0^\infty dr r^2 [E_h^2(r) + E_S^2(r) + E_P^2(r)], \quad (36)$$

where  $E_h(r)$ ,  $E_S(r)$  and  $E_P(r)$  are the equations of motion for the bubble wall profiles

$$E_h(r) \equiv \frac{d^2 h}{dr^2} + \frac{2}{r} \frac{dh}{dr} - \frac{\partial V_T^{eff}}{\partial h} = 0, \quad E_S(r) \equiv \frac{d^2 S}{dr^2} + \frac{2}{r} \frac{dS}{dr} - \frac{\partial V_T^{eff}}{\partial S} = 0,$$

$$E_P(r) \equiv \frac{d^2 P}{dr^2} + \frac{2}{r} \frac{dP}{dr} - \frac{\partial V_T^{eff}}{\partial P} = 0.$$

Note that the critical bubble obeys the following boundary conditions

$$(h(r), S(r), P(r)) \Big|_{r=\infty} = (0, S_s, P_s), \quad \left( \frac{dh}{dr}, \frac{dS}{dr}, \frac{dP}{dr} \right) \Big|_{r=0} = (0, 0, 0).$$

In Fig. 3 we show dependence of the critical scalar fields on the radial coordinate for the selected benchmark models at their critical temperatures. The corresponding values of the relevant physical parameters are shown in Table 2. All dimensionful parameters in Table 2 are in GeV. We observe considerable change in the values of the pseudoscalar field

	$v_S$	$v_P$	$\text{Im } \mu$	$T_c$	$v_c$	$S_c$	$P_c$	$S_s$	$P_s$	$S_3/T_c$
Setup 1	60	220	151.68	67.5	233.29	60.04	219.35	234.53	11.79	136.41
Setup 2	47.5	202.5	149.7	79.0	220.88	47.59	201.29	213.31	18.65	139.58

Table 2: Parameters for the first order EWPT in the split NMSSM.

$P$  in the broken and in the symmetric phases. This will be the source of  $CP$ -violation for generation of the baryon asymmetry during the EWPT.

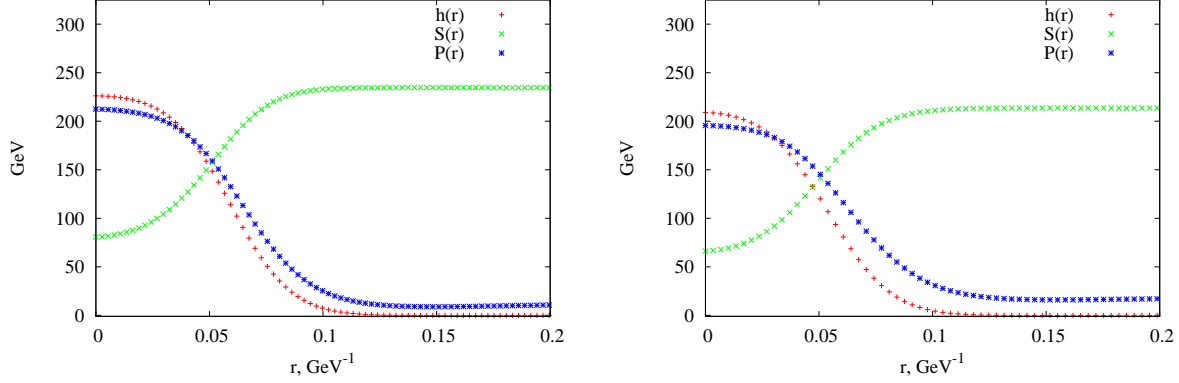


Figure 3: The critical bubble profile for the parameter set presented in Tables 1 and 2. Left and right panels correspond to Setups (1) and (2), respectively.

## 5 Baryon asymmetry

In this Section we discuss the baryon asymmetry created during the EWPT in the hot electroweak plasma. We will closely follow Ref. [34]. To the linear order in chemical potentials  $\mu_i$ , the particle asymmetry number density for  $i$ -th component of plasma reads as

$$n_i = \frac{1}{6} k_i \mu_i T^2, \quad (37)$$

where  $k_i$  equals 2 and 1 for massless boson and fermion degree of freedom, respectively. For nonrelativistic particle with mass  $m_i$ , the parameter  $k_i$  is suppressed by factor  $\exp(-m_i/T)$ .

It is convenient to introduce the following notations:  $n_T \equiv n_{t_R}$ ,  $n_Q \equiv n_{t_L} + n_{b_L}$ ,  $n_B \equiv n_{b_R}$ ,  $n_{H'} = n_{H^+} + n_{H^0}$ ,  $n_{\tilde{H}_u} = n_{\tilde{H}_u^+} + n_{\tilde{H}_u^0}$ ,  $n_{\tilde{H}_d} = n_{\tilde{H}_d^+} + n_{\tilde{H}_d^0}$ , and similar ones for the corresponding chemical potentials. We emphasize that the densities  $n_i$  are local quantities and through the parameters in (37) depend on  $z + v_w t$ , where  $z$  is the coordinate perpendicular to the bubble wall, and  $v_w$  is the wall velocity. The baryon number conservation implies the following relation  $n_B + n_T + n_Q = 0$ .

In diffusion equations we take into account scattering processes involving the top Yukawa coupling  $y_t \bar{Q} t H$  with the rate  $\Gamma_Y$ , strong sphaleron transitions with the rate  $\Gamma_{ss} = 6\kappa_{ss} \frac{8}{3} \alpha_s^4 T$ , Higgs boson self interactions with rate  $\Gamma_H$ . We also include the rate  $\Gamma_m$  for top quark mass interactions and the rates  $\Gamma_{u,d}$  for Higgs-gaugino-higgsino interactions. In addition, we take into account the Higgsino flipping interaction  $\tilde{\mu} \tilde{H}_u \epsilon \tilde{H}_d$  which has the rate  $\Gamma_\mu$ .

Following Ref. [34], we write down the set of diffusion equations in the large  $\tan \beta$

limit<sup>5</sup>

$$v_w n'_Q = D_q n''_Q - \Gamma_Y \left[ \frac{n_Q}{k_Q} - \frac{n_T}{k_T} - \frac{n_H + n_h}{k_{H_1}} \right] - \Gamma_m \left[ \frac{n_Q}{k_Q} - \frac{n_T}{k_T} \right] \quad (38)$$

$$-6\Gamma_{ss} \left[ 2\frac{n_Q}{k_Q} - \frac{n_T}{k_T} + 9\frac{n_Q + n_T}{k_B} \right],$$

$$v_w n'_T = D_q n''_T + \Gamma_Y \left[ \frac{n_Q}{k_Q} - \frac{n_T}{k_T} - \frac{n_H + n_h}{k_{H_1}} \right] + \Gamma_m \left[ \frac{n_Q}{k_Q} - \frac{n_T}{k_T} \right] \quad (39)$$

$$+ 3\Gamma_{ss} \left[ 2\frac{n_Q}{k_Q} - \frac{n_T}{k_T} + 9\frac{n_Q + n_T}{k_B} \right],$$

$$v_w n'_H = D_h n''_H + \Gamma_Y \left[ \frac{n_Q}{k_Q} - \frac{n_T}{k_T} - \frac{n_H + n_h}{k_{H_1}} \right] - \Gamma_H \frac{n_H + n_h}{k_{H_1}} + S_u - S_d, \quad (40)$$

$$v_w n'_h = D_h n''_h + \Gamma_Y \left[ \frac{n_Q}{k_Q} - \frac{n_T}{k_T} - \frac{n_H + n_h}{k_{H_1}} \right] - 2 \left[ \frac{n_H}{k_H} + \frac{n_h}{k_h} \right] \Gamma_{\tilde{\mu}} \quad (41)$$

$$- \frac{n_H + n_h}{k_{H_1}} \Gamma_H + S_u + S_d,$$

which are written for combinations of the Higgs bosons and higgsino densities  $n_h = n_{H'} + n_{\tilde{H}_u} + n_{\tilde{H}_d}$  and  $n_H = n_{H'} + n_{\tilde{H}_u} - n_{\tilde{H}_d}$ . Here the prime and double prime denote the first and second derivatives with respect to variable  $z$ . In equations (38)-(41), we set  $\Gamma_{\tilde{\mu}} \equiv \Gamma_{\mu} + \Gamma_d$ , statistical factors are defined by

$$k_{H_1} = 2(k_{H'} + k_{\tilde{H}_u}), \quad k_H = \frac{2(k_{H'} + k_{\tilde{H}_u})k_{\tilde{H}_d}}{k_{\tilde{H}_d} - k_{H'} - k_{\tilde{H}_u}}, \quad k_h = \frac{2(k_{H'} + k_{\tilde{H}_u})k_{\tilde{H}_d}}{k_{\tilde{H}_d} + k_{H'} + k_{\tilde{H}_u}},$$

and  $CP$ -violating sources  $S_u$  and  $S_d$  are discussed below in due course. We assume that the diffusion coefficients  $D_q$  are the same for all quarks, and  $D_h$  are the same for all Higgs bosons and higgsinos. Using the approach advocated in Ref. [35], we eliminate  $n_T$  and  $n_Q$  from Eqs. (38)-(41) by substituting the relations

$$n_T = -\frac{k_T(2k_B + 9k_Q)}{k_{H_1}(9k_Q + 9k_T + k_B)}(n_H + n_h), \quad n_Q = \frac{k_Q(9k_T - k_B)}{k_{H_1}(9k_Q + 9k_T + k_B)}(n_H + n_h), \quad (42)$$

which follow from the assumption that both the top Yukawa interactions and the strong sphalerons are in equilibrium. The resulting equations for  $n_H$  and  $n_h$  are collected in Ref. [25]. It follows from (42), that the left-handed fermion density can be recasted in the following form

$$n_{Left} = n_{Q_1} + n_{Q_2} + n_{Q_3} = 4n_T + 5n_Q$$

$$= \frac{5k_Q k_B + 8k_T k_B - 9k_Q k_T}{k_{H_1}(k_B + 9k_Q + 9k_T)}(n_h + n_H) \equiv A_t \cdot (n_h + n_H). \quad (43)$$

The statistical factors are  $k_Q = 6$ ,  $k_T = 3$ ,  $k_B = 3$ ,  $k_{H'} = 4$ ,  $k_{\tilde{H}_u} = k_{\tilde{H}_d} = 2$  and hence the constant  $A_t$  is equal to zero [36]. It was shown in Ref. [37], that one-loop corrections to statistical coefficients  $k_i$  give non-zero value of  $A_t$ , namely

$$A_t = \frac{3}{2k_{H_1}} \left( -\frac{3y_t^2}{8\pi^2} \right). \quad (44)$$

---

<sup>5</sup>At  $\tan \beta \sim 1$  the generated baryon asymmetry turns out to be suppressed, see Ref. [25] for details.

The baryon asymmetry obeys the following equation [35]

$$v_w n'_B(z) = -\theta(-z) [n_F \Gamma_{ws} n_{Left}(z) + \mathcal{R} n_B], \quad (45)$$

where  $\Gamma_{ws} = 6\kappa_{ws}\alpha_w^5 T$  is the weak sphaleron rate with  $\kappa_{ws} = 20 \pm 2$  [38]. The relaxation coefficient  $\mathcal{R}$  is given by [39]  $\mathcal{R} = \frac{5}{4}n_F\Gamma_{ws}$ , and  $n_F$  is the number of generations,  $n_F = 3$ . Here the domain  $z < 0$  corresponds to the symmetric phase. The solution to Eq. (45) reads

$$n_B = -\frac{n_F \Gamma_{ws}}{v_w} \int_{-\infty}^0 dz n_{Left}(z) e^{z\mathcal{R}/v_w}. \quad (46)$$

In the split NMSSM,  $CP$ -symmetry gets violated spontaneously while the bubble walls expand in the hot plasma. Indeed, the main source of  $CP$ -violation is associated with the complex chargino mass matrix

$$M_{ch} = \begin{pmatrix} M_2 & \frac{1}{\sqrt{2}}\tilde{g}_u h(z) \\ \frac{1}{\sqrt{2}}\tilde{g}_d h(z) & \tilde{\mu}(z) \end{pmatrix}, \quad (47)$$

where we define the spatially-dependent effective higgsino mass parameter as follows

$$\tilde{\mu}(z) = \mu + \kappa (S(z) + iP(z)) / \sqrt{2}. \quad (48)$$

In the above expressions,  $h(z)$ ,  $S(z)$  and  $P(z)$  are the kink approximations of the bubble walls [35]

$$h(z) = \frac{1}{2}v_c \left( 1 - \tanh \left[ \alpha \left( 1 - \frac{2z}{L_w} \right) \right] \right), \quad (49)$$

$$\begin{pmatrix} S(z) \\ P(z) \end{pmatrix} = \begin{pmatrix} S_c \\ P_c \end{pmatrix} - \frac{1}{2} \begin{pmatrix} \Delta S \\ \Delta P \end{pmatrix} \left( 1 + \tanh \left[ \alpha \left( 1 - \frac{2z}{L_w} \right) \right] \right), \quad (50)$$

here  $v_c$ ,  $S_c$  and  $P_c$  are the critical values of the scalar fields (see, e.g. Table 2),  $\Delta S \equiv S_c - S_s$  and  $\Delta P \equiv P_c - P_s$ . We set velocity of the bubble wall equal to  $v_w = 0.1$ , the coefficient  $\alpha$  is taken to be  $3/2$ . The bubble wall width  $L_w$  may be chosen in the range  $5/T_c < L_w < 30/T_c$  consistent with the special study [11] and the WKB thick-wall restriction,  $L_w T_c > 1$ .

Following Ref. [11], we define the rates in Eqs. (38-41) as

$$\Gamma_H = 0.0036T\theta(z - 0.5L_w), \quad \Gamma_m = 0.05T\theta(z - 0.5L_w), \quad \Gamma_{\tilde{\mu}} = 0.1T,$$

with  $\theta$  being the step function, and choose the diffusion coefficients in the form  $D_q = 6/T$  and  $D_h = 110/T$ .

We use the expressions for  $CP$ -violating sources  $S_d$  and  $S_u$  from Ref. [11] and numerically solve the set of diffusion equations for  $n_h(z)$  and  $n_H(z)$ . Then, we calculate the asymmetry of left fermions using Eq. (43) and by evaluating the integral (46) we obtain the baryon asymmetry generated during EWPT.

Let us consider the baryon-to-entropy ratio  $\Delta_B = n_B/s$  with the entropy density

$$s = 2\pi^2 g_{eff} T^3 / 45,$$

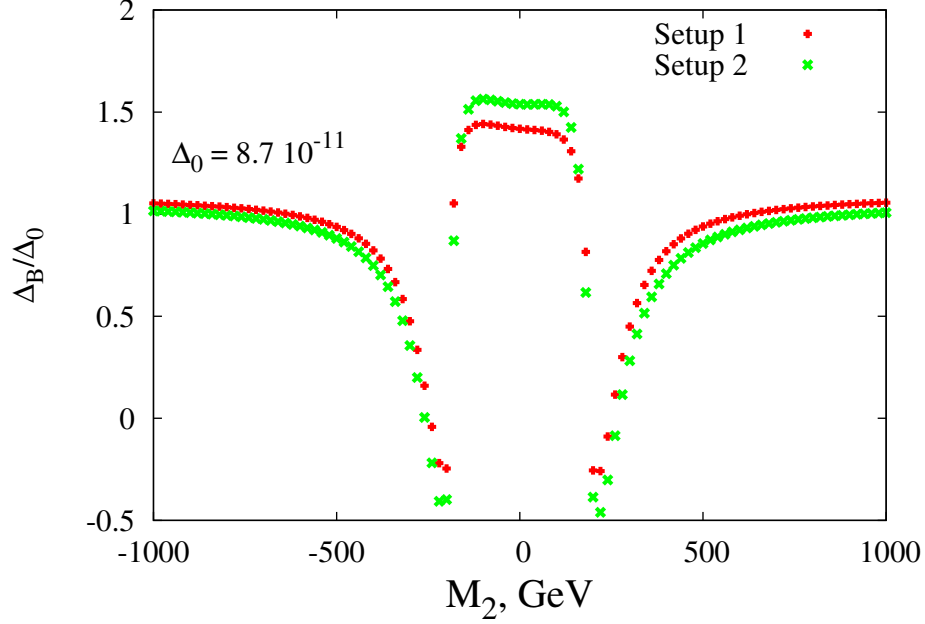


Figure 4: Plot of  $\Delta_B/\Delta_0$  versus gaugino mass parameter  $M_2$  for the parameter sets presented in Tables 2 and 1.

where  $g_{eff}$  is the effective number of relativistic degrees of freedom at  $T_c$ . In Fig. 4 we show dependence of the baryon asymmetry  $\Delta_B/\Delta_0$  on gaugino mass  $M_2$  for different values of the wall thickness: namely, we take  $L_w = 7/T_c$  and  $L_w = 5/T_c$  for Setup 1 and 2, respectively. The value  $\Delta_0 = 8.3 \times 10^{-11}$  corresponds to  $n_B/n_\gamma = 6.2 \times 10^{-10}$  consistent with present measurements (1).

It follows from Fig. 4 that baryon asymmetry  $\Delta_B$  is of order  $\Delta_0$  for large  $M_2 \gtrsim 1$  TeV. In this case, the heaviest chargino  $\chi_2^+$  (wino-like) decouples from the plasma,  $|m_{\chi_2^+}| \simeq M_2$ , and the lightest chargino (higgsino-like) acquires the mass  $|m_{\chi_1^+}|$ , which is determined by the effective  $\tilde{\mu}(z)$ -parameter in (48). Thus, the baryon asymmetry is generated due to the spontaneous  $CP$ -violation in the broken (and symmetric) phase. Detailed calculation of  $CP$ -violating sources [11] reveals that  $S_u$  and  $S_d$  gain contributions which are proportional to the second derivative of  $Im \tilde{\mu}(z)$  with respect to  $z$  coordinate. This means that baryon asymmetry  $\Delta_B/\Delta_0$  is rather sensitive to the effective parameter,

$$Im(\tilde{\mu}'') \sim \kappa \Delta P / L_w^2.$$

In our numerical analysis, we tune the wall thickness  $L_w$  to obtain  $\Delta_B/\Delta_0 \sim 1$  as  $M_2 \rightarrow \infty$ . At the same time from the very beginning we choose sufficiently large coupling  $\kappa$  (by taking large  $\lambda$ ) and pseudoscalar VEV gradient  $\Delta P = P_c - P_s$  and large value of  $\tan \beta$ . These features select models which are interesting for the realistic electroweak baryogenesis. As we will see in Section 6 the latter condition is also preferred by present electron's EDM constraints. From Fig. 2 we see that large values of  $\lambda$  and  $\tan \beta$  require moderate value of the splitting scale  $M_S$ , which hardly can be larger than 12 – 15 TeV.

In our analysis we can evaluate the baryon asymmetry in the limit  $n_h \gg n_H$ , following the approach, presented in Ref. [40]. In this approximation the set of diffusion equations (38-41) reduces to a single equation on  $n_h$ , and baryon asymmetry ratio,  $\Delta_B/\Delta_0$ , can be estimated analytically. In the limit when heaviest chargino decoupled,  $m_{\chi_2^+} \approx M_2 \approx 1$  TeV, one finds

$$\frac{\Delta_B}{\Delta_0} \approx 5.5 \cdot 10^2 \left( \frac{m_{\chi_1^+}}{T_c} \right)^2 \exp \left( -\frac{m_{\chi_1^+}}{T_c} \right) \frac{1}{(L_w T_c)^2}. \quad (51)$$

For  $L_w = 5/T_c$ ,  $T_c = 80$  GeV and  $m_{\chi_1^+} = 239$  GeV this yields  $\Delta_B/\Delta_0 \approx 10$ . An order-of-magnitude discrepancy between the numerical,  $\Delta_B/\Delta_0 \approx 1$ , and analytic results (51), is due to the approximations which have been made for solving equation for  $n_h$  in the analytically approach. Let us note that here we estimate baryon asymmetry originated from chargino sector only.  $CP$ -violating sources from neutralino sector can change the calculated value of the asymmetry by a factor of order one.

## 6 EDM constraints and light chargino phenomenology

In this Section we address some phenomenological implementation of the results discussed above. To begin with, we emphasize that current constraints on electric dipole moments of the electron and neutron provide strong limits for  $CP$ -violating physics in the split NMSSM. There are three relevant contributions to the EDM of electron or light quark [41, 44],

$$d_f = d_f^{H\gamma} + d_f^{HZ} + d_f^{WW},$$

where  $d_f^{H\gamma}$ ,  $d_f^{HZ}$  and  $d_f^{WW}$  are the partial EDMs of fermion (lepton or quark), related to the exchange of  $H\gamma$ ,  $HZ$  and  $W^+W^-$  bosons, respectively. General expressions for the electron's EDM  $d_e$  and neutron's EDM  $d_n$  were derived in Ref. [44]. The values of  $d_e$  and  $d_n$  depend on chargino,  $m_{\chi_i^+}$  ( $i = 1, 2$ ), and neutralino,  $m_{\chi_j^0}$  ( $j = 1, 5$ ), masses as well as their mixing matrices<sup>6</sup>.

The most stringent upper limit on EDM of the electron,  $|d_e/e| < 8.7 \times 10^{-29}$  cm at 90% CL, was obtained by ACME collaboration [45]. The current bound on neutron's EDM is  $|d_n/e| < 3.0 \times 10^{-26}$  cm at 90 % CL [46]. In order to perform the numerical analysis for EDMs, we randomly scan over the following parameter space  $0 < M_1, M_2 < 1000$  GeV. In Fig. 5 we show dependence of  $|d_e/e|$  on the lightest chargino mass  $m_{\chi_1^+}$ . One can see from the left panel of Fig. 5 that chargino masses in the ranges  $225 \text{ GeV} < m_{\chi_1^+} < 239 \text{ GeV}$  and  $220 \text{ GeV} < m_{\chi_1^+} < 235 \text{ GeV}$  are allowed for the Setup 1 and 2, respectively. We check that all these points correspond to large (about 1 TeV) values of  $M_2$  and hence allow for correct value of the BAU. The numerical results for neutron EDM are shown on right panel of Fig. 5. One can see from Fig. 5 that predictions for the neutron EDMs satisfy the current experimental bound in all selected models. For the Setup 1 we present

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<sup>6</sup>We recall that neutralino state  $\chi_j^0$  in split NMSSM is determined by the mixing of neutral bino  $\tilde{B}^0$ , wino  $\tilde{W}^0$ , higgsino  $\tilde{H}_{u,d}^0$  and singlino  $\tilde{n}$  states.

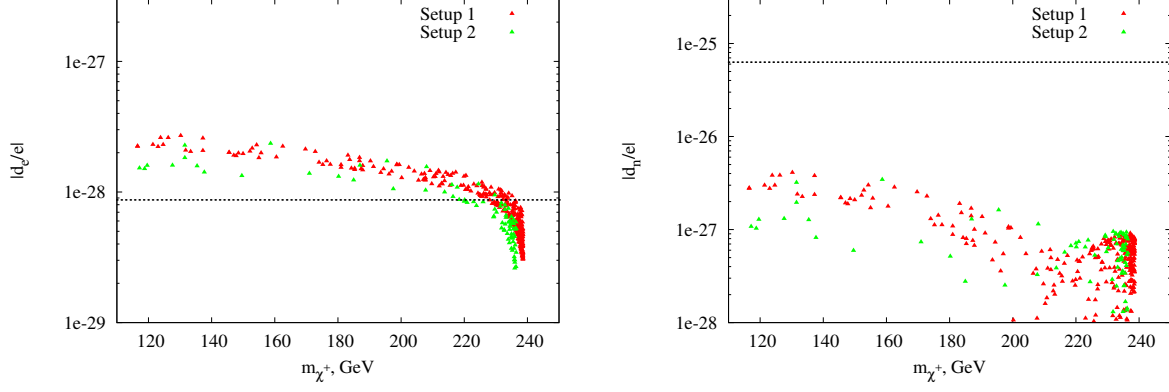


Figure 5: Left panel: the EDM of electron versus the lightest chargino mass  $m_{\chi_1^+}$ . Dotted lines represent the current experimental bound  $|d_e/e| < 8.7 \times 10^{-29}$  cm. Right panel: the neutron’s EDM with upper limit  $|d_n/e| < 3.0 \times 10^{-26}$  cm. The relevant couplings,  $\mu$ -terms and both singlet VEVs  $v_S$  and  $v_P$  at  $T = 0$  are given in Tables 1 and 2.

an examples of chargino and neutralino mass spectra which are consistent with the EDM bounds in Fig. 5 for  $M_1 = 300$  GeV and  $M_2 = 1$  TeV

- $m_{\chi_1^+} = 238.4$  GeV,  $m_{\chi_2^+} = 1006.8$  GeV,
- $m_{\chi_1^0} = 133.9$  GeV,  $m_{\chi_2^0} = 220.5$  GeV,  $m_{\chi_3^0} = 268.0$  GeV,  $m_{\chi_4^0} = 341.9$  GeV,  $m_{\chi_5^0} = 1006.9$  GeV.

We find in this case, that LSP is singlino-like state with the mass  $m_{\chi_1^0} = 133.9$  GeV. The dominant decay channel of the lightest chargino is  $\chi_1^+ \rightarrow \chi_1^0 W^+$ , which can be used to test split NMSSM model. In our analysis, we checked that the models satisfying EDM bounds are in agreement with the present CMS [47] and ATLAS [48] limits on chargino-neutralino production at LHC without light sleptons. Therefore, the split NMSSM is a phenomenologically viable and cosmologically attractive model which can be probed at the LHC run with  $pp$  collision energy of 13 TeV (and 14 TeV).

## 7 Conclusion

In this paper we revisit scenario of non-minimal split supersymmetry with possibility of realistic electroweak baryogenesis. It is a realization of split supersymmetry in the framework of NMSSM and contains at the electroweak scale, apart from minimal split supersymmetry particle content, singlet scalar and pseudoscalar states. We observed that within the phenomenologically allowed domain of the parameter space with the mass of the Higgs boson equal to 125 GeV it is possible to find particular models in which the strongly first order electroweak phase transition can be realized and moreover the needed amount of the baryon asymmetry of the Universe is generated. These models predict existence of light chargino state required for successful baryogenesis. We also find relatively light LSP with large admixture of singlino like state. Therefore, it can be



considered as a potential dark matter candidate as suggested in Ref. [25]. Predictions for the electric dipole moment of electron in these models are found to be about or somewhat larger than  $2 - 3 \cdot 10^{-29} e \text{ cm}$  which is only by factor 3-4 smaller than the current upper limit on this quantity. This makes the searches for EDMs a promising tool to probe the split NMSSM.

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## A One loop corrections to Higgs mass in split NMSSM

In this appendix we calculate one-loop RG corrections to the mass of Higgs boson in split NMSSM scenario following Refs. [26, 50]. In particular, in Ref. [50] the radiative corrections to the Higgs mass were calculated in the NMSSM in Ref. [50], while they were derived explicitly in split MSSM in Ref. [26]. However, split MSSM computations [26] can be straightforwardly extended to the split NMSSM case by taking into account the radiative corrections from scalars, charginos and neutralinos,

$$(m_h^{pole})^2 = (m_h^{tree})^2(\bar{\mu}) + \delta_h^{SM}(\bar{\mu}) + \delta_h^{(S,P)}(\bar{\mu}) + \delta_h^{(C,N)}(\bar{\mu}), \quad (52)$$

where  $(m_h^{tree})^2 = \tilde{\lambda}(\bar{\mu})v^2$  is the three level Higgs boson mass at  $\bar{\mu}$  scale (dimensional renormalization scale in  $\overline{\text{MS}}$  scheme); the remnant one-loop corrections in (52) are defined below in Sections A.1 and A.2. We use the experimental value of the Higgs pole mass (52) to plot the figures for the allowed region of split NMSSM parameters in the main text.

### A.1 Tree level potential of scalar sector in the broken phase

Applying the general results of Ref. [50] we rewrite (9) in the broken phase,

$$H = (\phi_1 + v)/\sqrt{2}, \quad N = (\phi_2 + v_S + i(\phi_3 + v_P))/\sqrt{2}, \quad (53)$$

where we denote perturbations of the scalar fields about the vacuum as  $(\phi_1, \phi_2, \phi_3) = (h, S, P)$ . Then, substituting (53) into (9) and using minimization conditions (88-89) at the tree level, one can obtain

$$\mathcal{L}_V \supset - \sum_{ijkl} \lambda_{\phi_i \phi_j \phi_k \phi_l} \phi_i \phi_j \phi_k \phi_l - \sum_{ijk} \lambda_{\phi_i \phi_j \phi_k} \phi_i \phi_j \phi_k - \sum_{ij} \frac{1}{2} m_{\phi_i \phi_j}^2 \phi_i \phi_j. \quad (54)$$

The quartic and trilinear couplings which are relevant for the calculation of the Higgs boson self energy and tadpoles in the scalar sector of the split NMSSM can be written as

$$\lambda_{\phi_1 \phi_1 \phi_1 \phi_1} = \frac{1}{8} \tilde{\lambda}, \quad \lambda_{\phi_1 \phi_1 \phi_2 \phi_2} = \frac{1}{12} (\kappa_1 + \kappa_2), \quad \lambda_{\phi_1 \phi_1 \phi_3 \phi_3} = \frac{1}{12} (\kappa_1 - \kappa_2), \quad (55)$$

$$\lambda_{\phi_1 \phi_1 \phi_1} = \frac{1}{2} \tilde{\lambda} v, \quad \lambda_{\phi_1 \phi_2 \phi_2} = \frac{1}{3} (\kappa_1 + \kappa_2) v, \quad \lambda_{\phi_1 \phi_3 \phi_3} = \frac{1}{3} (\kappa_1 - \kappa_2) v, \quad (56)$$

$$\lambda_{\phi_1\phi_3\phi_2} = \lambda_{\phi_1\phi_2\phi_3} = 0. \quad (57)$$

The parameters of the scalar squared mass matrix read

$$m_{\phi_3\phi_3}^2 = (\kappa_1 - \kappa_2)v^2 + \lambda_N(3v_P^2 + v_S^2) - \tilde{\lambda}(v_P^2 + v_S^2), \quad (58)$$

$$m_{\phi_2\phi_3}^2 = m_{\phi_3\phi_2}^2 = -\sqrt{2}\tilde{A}_k v_P + 2\lambda_N v_P v_S. \quad (59)$$

$$m_{\phi_1\phi_1}^2 = \tilde{\lambda}v^2, \quad m_{\phi_2\phi_2}^2 = (\kappa_1 + \kappa_2)v^2 + \lambda_N(v_P^2 + 3v_S^2) + \left(-\tilde{\lambda} + \tilde{A}_k/(\sqrt{2}v_S)\right)(v_P^2 + v_S^2), \quad (60)$$

$$m_{\phi_1\phi_3}^2 = m_{\phi_3\phi_1}^2 = m_{\phi_1\phi_2}^2 = m_{\phi_2\phi_1}^2 = 0. \quad (61)$$

One should diagonalize its  $2 \times 2$  submatrix for the singlets  $m_{\phi_i\phi_j}^2$ , with  $i, j = 2, 3$ , since off-diagonal mixings of  $\phi_2$  and  $\phi_3$  with the Higgs field  $\phi_1$  are set to be zero (61) (see also discussion before Eq. (16)). We denote the singlet eigenstates by  $h_i$  and diagonalize  $m_{\phi_i\phi_j}^2$  by an orthogonal matrix  $R_{ij}$ , such that

$$h_i = R_{ij}\phi_j. \quad (62)$$

The couplings that enter the calculation of the Higgs boson mass radiative corrections can be expressed as

$$\lambda_{\phi_i\phi_j h_k h_l} = 6 R_{ka} R_{lb} \lambda_{\phi_i\phi_j\phi_a\phi_b}, \quad \lambda_{\phi_i h_k h_l} = 3 R_{ka} R_{lb} \lambda_{\phi_i\phi_a\phi_b}. \quad (63)$$

Following the prescription of Ref. [50] we write down one-loop contribution of the scalar singlets to the Higgs boson mass <sup>7</sup>

$$\delta_h^{(S,P)}(m_h, \bar{\mu}) = \frac{1}{v} T_h^{(S,P)}(\bar{\mu}) - \Pi_h^{(S,P)}(m_h, \bar{\mu}), \quad (64)$$

where the Higgs boson self energy is

$$16\pi^2 \Pi_h^{(S,P)}(p^2, \bar{\mu}) = \sum_{k=2,3} 2\lambda_{\phi_1\phi_1 h_k h_k} A_0(m_{h_k}) + \sum_{k,l=2,3} 2\lambda_{\phi_1 h_k h_l} \lambda_{\phi_1 h_k h_l} B_0(p, m_{h_k}, m_{h_l}), \quad (65)$$

and the tadpole contributions are

$$16\pi^2 T_h^{(S,P)}(\bar{\mu}) = \sum_{k=2,3} \lambda_{\phi_1 h_k h_k} A_0(m_{h_k}). \quad (66)$$

The loop functions  $A_0(m)$  and  $B_0(p, m_1, m_2)$  depend on the renormalization scale  $\bar{\mu}$  and can be written in the form

$$A_0(m) = m^2 \left( C_{UV} + 1 - \ln \frac{m^2}{\bar{\mu}^2} \right), \quad B_0(p, m_1, m_2) = C_{UV} - \ln \frac{p^2}{\bar{\mu}^2} - f_B(x_+) - f_B(x_-), \quad (67)$$

where  $C_{UV} = 1/\epsilon - \gamma_E + \ln 4\pi$ ,  $f_B(x) = \ln(1-x) - x \ln(1-x^{-1}) - 1$  with

$$x_{\pm} = \frac{s \pm \sqrt{s^2 - 4p^2(m_1^2 - i\epsilon)}}{2p^2}, \quad s = p^2 - m_2^2 + m_1^2. \quad (68)$$

A simplified formula for  $B_0(p^2, m_1, m_2)$  at  $p^2 = 0$  read [51],

$$B_0(0, m_1, m_2) = -\ln \frac{M^2}{\bar{\mu}^2} + 1 + \frac{m^2}{m^2 - M^2} \ln \frac{M^2}{m^2}, \quad (69)$$

where  $M = \max(m_1, m_2)$  and  $m = \min(m_1, m_2)$ .

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<sup>7</sup>Here only scalars  $\phi_2$  and  $\phi_3$  are taken into account; all signs and prefactors correspond to notations from Ref. [50].

## A.2 Chargino-neutralino sector of split NMSSM

The Lagrangian of interest for chargino/neutralino sector is

$$\begin{aligned}
-\mathcal{L}_{int}^{(C,N)} = & -\frac{1}{2}h \bar{\chi}_i^0 (R_{ij}^{N*} P_L + R_{ij}^N P_R) \chi_j^0 + \\
& + \left( g \bar{\chi}_i^+ \gamma^\mu (C_{ij}^R P_R + C_{ij}^L P_L) \chi_j^0 W_\mu^+ + \frac{1}{\sqrt{2}} \bar{\chi}_i^+ (R_{ij}^C P_R + L_{ij}^C P_L) \chi_j^+ h + h.c. \right),
\end{aligned} \tag{70}$$

where

$$R_{ij}^N = (\tilde{g}_u N_{i2} - \tilde{g}_u' N_{i1}) N_{j4} - (\tilde{g}_d N_{i2} - \tilde{g}_d' N_{i1}) N_{j3} + \sqrt{2}(\lambda_u N_{i4} - \lambda_d N_{i3}) N_{j5} \tag{71}$$

$$R_{(ij)}^N = \frac{1}{2}(R_{ij}^N + R_{ji}^N), \quad R_{ij}^C = (L_{ji}^C)^* = \tilde{g}_u^* V_{i2} U_{j1} + \tilde{g}_d^* V_{i1} U_{j2}, \tag{72}$$

$$C_{ij}^L = N_{i2} V_{j1}^* - \frac{1}{\sqrt{2}} N_{i4} V_{j2}^*, \quad C_{ij}^R = N_{i2}^* U_{j1} + \frac{1}{\sqrt{2}} N_{i3}^* U_{j2}. \tag{73}$$

Following Ref. [26], let us consider the contribution of chargino and neutralino to the Higgs boson mass at one-loop level,

$$\delta_h^{(C,N)} = \Sigma_h^{(C,N)}(m_h, \bar{\mu}) + \frac{1}{v} T_h^{(C,N)}(\bar{\mu}) + \frac{\tilde{\lambda} v^2}{m_W^2} \Pi_{WW}^{(C,N)}(0, \bar{\mu}) \tag{74}$$

where  $T_h^{(C,N)} = T_h^{(C)} + T_h^{(N)}$  is the Higgs boson tadpole contribution which involves terms

$$16\pi^2 T^{(C)}(\bar{\mu}) = -2\sqrt{2} \sum_{i=1}^2 \text{Re} [R_{ii}^C M_i^C A_0(M_i^C)], \tag{75}$$

$$16\pi^2 T^{(N)}(\bar{\mu}) = 2 \sum_{i=1}^5 \text{Re} [R_{(ii)}^N M_i^N A_0(M_i^N)] \tag{76}$$

from chargino and and neutralino sector, respectively. The relevant self energies read  $\Sigma_h^{(C,N)} = \Sigma_h^{(C)} + \Sigma_h^{(N)}$ , where

$$16\pi^2 \Sigma_h^{(C)}(p^2, \bar{\mu}) = \sum_{i,j=1}^2 \left[ \frac{1}{2} (|L_{ij}^C|^2 + |R_{ij}^C|^2) \left( A_0(M_i^C) + A_0(M_j^C) + \right. \right. \tag{77}$$

$$\left. + ((M_i^C)^2 + (M_j^C)^2 - p^2) B_0(p^2, M_i^C, M_j^C) \right) + 2 \text{Re} M_i^C M_j^C R_{ij}^C (L_{ij}^C)^* B_0(p^2, M_i^C, M_j^C) \Big],$$

$$16\pi^2 \Sigma_h^{(N)}(p^2, \bar{\mu}) = \sum_{i,j=1}^5 \left[ |R_{(ij)}^N|^2 \left( A_0(M_i^N) + A_0(M_j^N) + \right. \right. \tag{78}$$

$$\left. + ((M_i^N)^2 + (M_j^N)^2 - p^2) B_0(p^2, M_i^N, M_j^N) \right) + 2 \text{Re} M_i^N M_j^N R_{(ij)}^N (R_{(ij)}^N)^* B_0(p^2, M_i^N, M_j^N) \Big].$$

The last term in Eq. (74) is the corrections from the contribution of chargino and neutralino into the  $W^\pm$  boson self-energy

$$\begin{aligned}
& 16\pi^2 \Pi_{WW}^{(C,N)}(0, \bar{\mu}) = \\
& = g^2 \sum_{i=1}^5 \sum_{j=1}^2 \left( (C_{ij}^L C_{ij}^{L*} + C_{ij}^R C_{ij}^{R*}) \left[ a^2 \left( \ln \frac{a^2}{\bar{\mu}^2} - \frac{1}{2} \right) + b^2 \left( \ln \frac{b^2}{\bar{\mu}^2} - \frac{1}{2} \right) + \frac{a^2 b^2}{a^2 - b^2} \ln \frac{a^2}{b^2} \right] + \right. \\
& \quad \left. + 2(C_{ij}^L C_{ij}^{R*} + C_{ij}^R C_{ij}^{L*}) \frac{ab}{a^2 - b^2} \left[ -a^2 \left( \ln \frac{a^2}{\bar{\mu}^2} - 1 \right) + b^2 \left( \ln \frac{b^2}{\bar{\mu}^2} - 1 \right) \right] \right), \quad (79)
\end{aligned}$$

where  $a = M_j^C$  and  $b = M_i^N$  are the mass eigenstates of chargino and neutralino, respectively. For the explicit calculation of Higgs mass (52), one should set  $C_{UV} = 0$  in (64) and (74).

### A.3 One-loop correction to Yukawa coupling of top quark

The mass of the Higgs boson at one-loop level is quite sensitive to the Yukawa coupling of top quark,  $y_t$ . Hence, it is important to include the RG effects and threshold corrections from top quark sector for explicit analysis of one-loop corrections to the Higgs boson mass in the split NMSSM. Here we briefly summarize the results of [26] concerning corrections related to  $y_t$ . The top quark Yukawa coupling at the scale  $\bar{\mu}$  can be extracted from its pole mass  $M_t = 173.2 \pm 0.9$  GeV [49],

$$y_t(\bar{\mu}) = \sqrt{2} \frac{M_t}{v} (1 + \delta_t(\bar{\mu})), \quad (80)$$

where the threshold correction  $\delta_t(\bar{\mu})$  is the sum of the QCD, EW and split NMSSM terms

$$\delta_t(\bar{\mu}) = \delta_t^{QCD}(\bar{\mu}) + \delta_t^{EW}(\bar{\mu}) + \delta_t^{(C,N)}(\bar{\mu}). \quad (81)$$

Explicit 3-loop calculation of  $\delta_t^{QCD}(\bar{\mu})$  was performed by [52] and at  $\bar{\mu} = M_t$  it yields

$$\delta_t^{QCD}(\bar{\mu} = M_t) = -\frac{4}{3} \left( \frac{\alpha_3(M_t)}{\pi} \right) - 9.1 \left( \frac{\alpha_3(M_t)}{\pi} \right)^2 - 80 \left( \frac{\alpha_3(M_t)}{\pi} \right)^3 \approx -0.060. \quad (82)$$

The contribution of the EW term  $\delta_t^{EW}$  is negligible [26],  $|\delta_t^{EW}| < 0.001$ . The term  $\delta_t^{(C,N)}$  from chargino and neutralino in split NMSSM is given through the relation

$$\delta_t(\bar{\mu}) = -\frac{\Pi_{WW}^{(C,N)}(0, \bar{\mu})}{2M_W^2}, \quad (83)$$

where  $\Pi_{WW}^{(C,N)}(0, \bar{\mu})$  is defined by Eq. (79).

## B Minimization of the effective potential

Here we present minimization conditions for the scalar potential of the model which allow us to express the soft parameters  $m^2$ ,  $\tilde{m}^2$  and  $\tilde{m}_N^2$  via the expectation values  $v$ ,  $v_S$  and  $v_P$  (cf. Eq. (9)). To do that, let us consider one-loop effective potential at zero temperature

$$V_{T=0}^{eff} = V_{tree} + V^{(1)}, \quad (84)$$

where  $V_{tree}$  is the tree level potential,  $V_{tree} \equiv -\mathcal{L}_V$ , and  $V^{(1)}$  is the one-loop contribution of fermions, gauge bosons and scalars to the effective potential. In the  $\overline{\text{DR}}$  scheme,  $V^{(1)}$  has the form [53]

$$V^{(1)} = \sum_i (\pm) \frac{n_i m_i^4}{64\pi^2} \left( \ln \frac{m_i^2}{q^2} - \frac{3}{2} \right), \quad (85)$$

here  $(+)$  is for bosons and  $(-)$  is for fermions and sum runs over all particles which have field-dependent mass  $m_i$  and  $n_i$  degrees of freedom. We choose the renormalization scale  $q$  at 100 GeV. In order to define the global minimum of the potential (84) at the fixed point  $(v, v_S, v_P)$ , we expand  $V_{tree}$  in the following way

$$V_{tree} = -\frac{m^2}{2}v^2 + \frac{\tilde{m}_N^2}{2}(v_S^2 + v_P^2) + \frac{\tilde{m}^2}{2}(v_S^2 - v_P^2) + V_{tree}^{>2}, \quad (86)$$

where  $V_{tree}^{>2}$  stands for cubic and quartic terms of the tree level potential (9). The vacuum  $(v, v_S, v_P)$  at zero temperature is determined by the stationary conditions

$$\frac{\partial}{\partial v} V_{T=0}^{eff}(v, v_S, v_P) = \frac{\partial}{\partial v_S} V_{T=0}^{eff}(v, v_S, v_P) = \frac{\partial}{\partial v_P} V_{T=0}^{eff}(v, v_S, v_P) = 0. \quad (87)$$

We emphasize that  $v$  is fine-tuned to be the vacuum expectation value of Higgs boson,  $v = 246$  GeV. It follows from Eq. (86) and Eq. (87), that squared masses  $m^2$ ,  $\tilde{m}_N^2$  and  $\tilde{m}^2$  can be redefined in the following form

$$m^2 = \frac{1}{v} \frac{\partial}{\partial v} (V_{tree}^{>2} + V^{(1)}), \quad \tilde{m}_N^2 = \frac{1}{2} \left( \frac{1}{v_S} \frac{\partial}{\partial v_S} + \frac{1}{v_P} \frac{\partial}{\partial v_P} \right) (V_{tree}^{>2} + V^{(1)}), \quad (88)$$

$$\tilde{m}^2 = \frac{1}{2} \left( \frac{1}{v_S} \frac{\partial}{\partial v_S} - \frac{1}{v_P} \frac{\partial}{\partial v_P} \right) (V_{tree}^{>2} + V^{(1)}). \quad (89)$$

At the tree level this yields

$$m^2 = \frac{1}{2} \tilde{\lambda} v^2 - (\kappa_1 - \kappa_2) v_P^2 - (\kappa_1 + \kappa_2) v_S^2, \quad \tilde{m}_N^2 = \frac{\tilde{A}_k}{\sqrt{2}} \left( \frac{v_S}{2} + \frac{v_P^2}{v_S} \right) - \frac{\lambda_N}{2} (v_S^2 + v_P^2),$$

$$\tilde{m}^2 = \frac{\tilde{A}_k}{\sqrt{2}} \left( \frac{v_P^2}{v_S} - \frac{3}{2} v_S \right) - \frac{1}{2} \lambda_N (v_P^2 - v_S^2).$$

Here we neglect contribution arising from  $V^{(1)}$ ,  $\xi$  and  $\eta$  terms. We stress that in this notation the singlet's VEVs,  $v_S$  and  $v_P$ , are the free dimensionful parameters of the split NMSSM. In our analysis we scan over these two VEVs to find the parameter space with successful baryogenesis and calculate the electron and neutron EDMs as obeying the present experimental constraints.

## References

- [1] C. L. Bennett *et al.* [WMAP Collaboration], *Astrophys. J. Suppl.* **208** (2013) 20 [arXiv:1212.5225 [astro-ph.CO]].

- [2] A. D. Sakharov, Pisma Zh. Eksp. Teor. Fiz. **5**, 32 (1967) [JETP Lett. **5** (1967 SOPUA,34,392-393.1991 UFNAA,161,61-64.1991) 24].
- [3] D. E. Morrissey and M. J. Ramsey-Musolf, New J. Phys. **14**, 125003 (2012) [arXiv:1206.2942 [hep-ph]].
- [4] T. Konstandin, Phys. Usp. **56**, 747 (2013) [Usp. Fiz. Nauk **183**, 785 (2013)] [arXiv:1302.6713 [hep-ph]].
- [5] K. Kajantie, M. Laine, K. Rummukainen and M. E. Shaposhnikov, Phys. Rev. Lett. **77**, 2887 (1996) [hep-ph/9605288].
- [6] F. Csikor, Z. Fodor and J. Heitger, Phys. Rev. Lett. **82**, 21 (1999) [hep-ph/9809291].
- [7] D. Bodeker, L. Fromme, S. J. Huber and M. Seniuch, JHEP **0502**, 026 (2005) [hep-ph/0412366].
- [8] L. Fromme, S. J. Huber and M. Seniuch, JHEP **0611**, 038 (2006) [hep-ph/0605242].
- [9] Y. Li, S. Profumo and M. Ramsey-Musolf, Phys. Lett. B **673**, 95 (2009) [arXiv:0811.1987 [hep-ph]].
- [10] J. Kozaczuk, S. Profumo, M. J. Ramsey-Musolf and C. L. Wainwright, Phys. Rev. D **86**, 096001 (2012) [arXiv:1206.4100 [hep-ph]].
- [11] S. J. Huber and M. G. Schmidt, Nucl. Phys. B **606**, 183 (2001) [arXiv:hep-ph/0003122].
- [12] S. J. Huber and M. G. Schmidt, arXiv:hep-ph/0011059.
- [13] S. J. Huber, T. Konstandin, T. Prokopec and M. G. Schmidt, Nucl. Phys. B **757**, 172 (2006) [hep-ph/0606298].
- [14] J. Kozaczuk, S. Profumo and C. L. Wainwright, Phys. Rev. D **87**, no. 7, 075011 (2013) [arXiv:1302.4781 [hep-ph]].
- [15] K. Cheung, T. J. Hou, J. S. Lee and E. Senaha, Phys. Lett. B **710**, 188 (2012) [arXiv:1201.3781 [hep-ph]].
- [16] J. L. Feng, Ann. Rev. Astron. Astrophys. **48**, 495 (2010) doi:10.1146/annurev-astro-082708-101659
- [17] S. P. Martin, Adv. Ser. Direct. High Energy Phys. **21**, 1 (2010), [hep-ph/9709356].
- [18] S. Chatrchyan *et al.* [CMS Collaboration], Phys. Lett. B **716**, 30 (2012) [arXiv:1207.7235 [hep-ex]].
- [19] G. Aad *et al.* [ATLAS Collaboration], Phys. Lett. B **716**, 1 (2012) [arXiv:1207.7214 [hep-ex]].

- [20] G. Aad *et al.* [ATLAS Collaboration], “Search for squarks and gluinos with the ATLAS detector in final states with jets and missing transverse momentum using 4.7 fb<sup>-1</sup> of  $\sqrt{s} = 7$  TeV proton-proton collision data,” Phys. Rev. D **87**, no. 1, 012008 (2013) [arXiv:1208.0949 [hep-ex]].
- [21] S. Chatrchyan *et al.* [CMS Collaboration], “Search for supersymmetry in hadronic final states using MT2 in  $pp$  collisions at  $\sqrt{s} = 7$  TeV,” JHEP **1210**, 018 (2012) [arXiv:1207.1798 [hep-ex]].
- [22] N. Arkani-Hamed and S. Dimopoulos, JHEP **0506** (2005) 073 [hep-th/0405159].
- [23] G. F. Giudice and A. Romanino, Nucl. Phys. B **699**, 65 (2004) [Erratum-ibid. B **706**, 65 (2005)] [arXiv:hep-ph/0406088].
- [24] K. A. Olive *et al.* [Particle Data Group Collaboration], Chin. Phys. C **38**, 090001 (2014).
- [25] S. V. Demidov and D. S. Gorbunov, JHEP **0702**, 055 (2007) [hep-ph/0612368].
- [26] M. Binger, Phys. Rev. D **73**, 095001 (2006) doi:10.1103/PhysRevD.73.095001 [hep-ph/0408240].
- [27] G. F. Giudice and A. Strumia, Nucl. Phys. B **858**, 63 (2012) doi:10.1016/j.nuclphysb.2012.01.001 [arXiv:1108.6077 [hep-ph]].
- [28] L. Dolan and R. Jackiw, Phys. Rev. D **9**, 3320 (1974).
- [29] G. D. Moore, Phys. Rev. D **59**, 014503 (1999)
- [30] H. H. Patel and M. J. Ramsey-Musolf, JHEP **1107** (2011) 029 doi:10.1007/JHEP07(2011)029 [arXiv:1101.4665 [hep-ph]].
- [31] G. W. Anderson and L. J. Hall, Phys. Rev. D **45**, 2685 (1992).
- [32] J. M. Moreno, M. Quiros and M. Seco, Nucl. Phys. B **526**, 489 (1998) [arXiv:hep-ph/9801272].
- [33] P. John, Phys. Lett. B **452**, 221 (1999) [arXiv:hep-ph/9810499]. S. J. Huber, P. John, M. Laine and M. G. Schmidt, Phys. Lett. B **475**, 104 (2000) [arXiv:hep-ph/9912278]. S. J. Huber, P. John and M. G. Schmidt, Eur. Phys. J. C **20**, 695 (2001) [arXiv:hep-ph/0101249].
- [34] P. Huet and A. E. Nelson, Phys. Rev. D **53**, 4578 (1996) [hep-ph/9506477].
- [35] M. Carena, J. M. Moreno, M. Quiros, M. Seco and C. E. M. Wagner, Nucl. Phys. B **599**, 158 (2001) [arXiv:hep-ph/0011055].
- [36] G. F. Giudice and M. E. Shaposhnikov, Phys. Lett. B **326**, 118 (1994) [arXiv:hep-ph/9311367].

- [37] M. Carena, A. Megevand, M. Quiros and C. E. M. Wagner, Nucl. Phys. B **716**, 319 (2005) [arXiv:hep-ph/0410352].
- [38] G. D. Moore and K. Rummukainen, Phys. Rev. D **61**, 105008 (2000) [arXiv:hep-ph/9906259].
- [39] J. M. Cline, M. Joyce and K. Kainulainen, Phys. Lett. B **417**, 79 (1998) [Erratum-ibid. B **448**, 321 (1999)] [arXiv:hep-ph/9708393]. J. M. Cline, M. Joyce and K. Kainulainen, JHEP **0007**, 018 (2000) [arXiv:hep-ph/0006119].
- [40] M. Carena, J. M. Moreno, M. Quiros, M. Seco and C. E. M. Wagner, Nucl. Phys. B **599**, 158 (2001) doi:10.1016/S0550-3213(01)00032-3 [hep-ph/0011055].
- [41] N. Arkani-Hamed, S. Dimopoulos, G. F. Giudice and A. Romanino, Nucl. Phys. B **709**, 3 (2005) [arXiv:hep-ph/0409232].
- [42] D. Chang, W. F. Chang and W. Y. Keung, Phys. Rev. D **71**, 076006 (2005) [arXiv:hep-ph/0503055].
- [43] N. G. Deshpande and J. Jiang, Phys. Lett. B **615**, 111 (2005) [arXiv:hep-ph/0503116].
- [44] G. F. Giudice and A. Romanino, Phys. Lett. B **634**, 307 (2006) [arXiv:hep-ph/0510197].
- [45] J. Baron *et al.* [ACME Collaboration], Science **343**, 269 (2014) [arXiv:1310.7534 [physics.atom-ph]].
- [46] C. A. Baker *et al.*, arXiv:hep-ex/0602020.
- [47] V. Khachatryan *et al.* [CMS Collaboration], Eur. Phys. J. C **74**, no. 9, 3036 (2014) [arXiv:1405.7570 [hep-ex]].
- [48] G. Aad *et al.* [ATLAS Collaboration], JHEP **1405**, 071 (2014) [arXiv:1403.5294 [hep-ex]].
- [49] [Tevatron Electroweak Working Group and CDF and D0 Collaborations], arXiv:1107.5255 [hep-ex].
- [50] G. Degrandi and P. Slavich, Nucl. Phys. B **825**, 119 (2010) doi:10.1016/j.nuclphysb.2009.09.018 [arXiv:0907.4682 [hep-ph]].
- [51] D. M. Pierce, J. A. Bagger, K. T. Matchev and R. j. Zhang, Nucl. Phys. B **491** (1997) 3 doi:10.1016/S0550-3213(96)00683-9 [hep-ph/9606211].
- [52] K. G. Chetyrkin and M. Steinhauser, Phys. Rev. Lett. **83** (1999) 4001 doi:10.1103/PhysRevLett.83.4001 [hep-ph/9907509].
- [53] S. R. Coleman and E. Weinberg, Phys. Rev. D **7**, 1888 (1973).